

Predictions for $\Xi_b^- \rightarrow \pi^- (D_s^-) \Xi_c^0(2790) (\Xi_c^0(2815))$ and $\Xi_b^- \rightarrow \bar{\nu}_l l \Xi_c^0(2790) (\Xi_c^0(2815))$

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(Dated: January 25, 2017)

We have performed the calculations for the nonleptonic $\Xi_b^- \rightarrow \pi^- \Xi_c^0(2790) (J = \frac{1}{2})$ and $\Xi_b^- \rightarrow \pi^- \Xi_c^0(2815) (J = \frac{3}{2})$ and the same reactions replacing the π^- by a D_s^- . At the same time we evaluate the semileptonic rates for $\Xi_b^- \rightarrow \bar{\nu}_l l \Xi_c^0(2790)$ and $\Xi_b^- \rightarrow \bar{\nu}_l l \Xi_c^0(2815)$. We look at the reactions from the perspective that the $\Xi_c^0(2790)$ and $\Xi_c^0(2815)$ resonances are dynamically generated from the pseudoscalar-baryon and vector-baryon interactions. We evaluate ratios of the rates of these reactions and make predictions that can be tested in future experiments. We also find that the results are rather sensitive to the coupling of the Ξ_c^* resonances to the $D^*\Sigma$ and $D^*\Lambda$ components.

PACS numbers:

I. INTRODUCTION

The introduction of chiral dynamics in the study of meson baryon interactions [1, 2] has allowed a rapid development in this field. A qualitative step forward was given by introducing unitarity in coupled channels, using the chiral Lagrangians as a source of the interaction [3–7]. In many cases the interaction is strong enough to generate bound states in some channels, which decay into the open states considered in the coupled channel formalism. The most renowned case is the one of the two $\Lambda(1405)$ states [5, 6, 8, 9]. The original works considered the interaction of pseudoscalar mesons with baryons, but the extension to vector mesons with baryons was soon done in Refs. [10, 11]. The extension to vector mesons finds its natural framework in the use of the local hidden gauge Lagrangians [12–14], which extend the chiral Lagrangians and accommodate vector mesons.

The mixing of pseudoscalar-baryon and vector-baryon channels in that framework was done in Ref. [15] in the light sector, and was extended to the charm sector in Refs. [16, 17]. An alternative approach to this mixing has been done in Ref. [18], where the chiral Weinberg-Tomozawa meson-baryon interaction was extended to the $SU(8)$ spin-flavour symmetry group.

One case where the relevance of the mixing is found is in the description of the $\Lambda_c(2595)(\frac{1}{2}^-)$ and $\Lambda_c(2625)(\frac{3}{2}^-)$. In early works on the subject the $\Lambda_c(2595)$ appeared basically as a DN molecule [19, 20], but both in Ref. [18] and Ref. [16] a coupling to the D^*N component was found with similar strength. On the other hand the $\Lambda_c(2625)$

appears from the D^*N interaction in S -wave.

Support for the relevance of the vector-baryon components in these states was recently found in Refs. [21, 22]. In Ref. [21] the decays $\Lambda_b \rightarrow \pi^- \Lambda_c(2595)$ and $\Lambda_b \rightarrow \pi^- \Lambda_c(2625)$ were studied and good agreement with experiment was found for the ratio of the two partial decay widths. The role of the D^*N was found very important, to the point that if the sign of the coupling of the D^*N to the $\Lambda_c(2595)$ was changed, the ratio of partial decay widths was in sheer disagreement with experiment. In Ref. [22] the semileptonic decay $\Lambda_b \rightarrow \bar{\nu}_l l \Lambda_c(2595)$ and $\Lambda_b \rightarrow \bar{\nu}_l l \Lambda_c(2625)$ were studied and the ratio of the partial decay widths was also found in agreement with experiment. Once again, reversing the sign of the D^*N coupling to the $\Lambda_c(2595)$ led to results incompatible with experiment.

In the present work we retake the ideas of Refs. [21, 22] and apply them to the study of the decays $\Xi_b^- \rightarrow \pi^- \Xi_c^0(2790)(\frac{1}{2}^-)$, $\Xi_b^- \rightarrow \pi^- \Xi_c^0(2815)(\frac{3}{2}^-)$, $\Xi_b^- \rightarrow D_s^- \Xi_c^0(2790)$, $\Xi_b^- \rightarrow D_s^- \Xi_c^0(2815)$, $\Xi_b^- \rightarrow \bar{\nu}_l l \Xi_c^0(2790)$, $\Xi_b^- \rightarrow \bar{\nu}_l l \Xi_c^0(2815)$. The $\Xi_c^0(2790)(\frac{1}{2}^-)$ and $\Xi_c^0(2815)(\frac{3}{2}^-)$ play an analogous role to the $\Lambda_c(2595)(\frac{1}{2}^-)$ and $\Lambda_c(2625)(\frac{3}{2}^-)$, substituting the u -quark by an s -quark. In Ref. [18] the couplings of the $\Xi_c^0(2790)$ and $\Xi_c^0(2815)$ to the different coupled channels were evaluated for both pseudoscalar-baryon and vector-baryon components, in particular the $D\Lambda$, $D^*\Lambda$, $D\Sigma$, $D^*\Sigma$ which will be those needed in the decays mentioned above. We will adapt the formalism developed in Refs. [21, 22] to the present case and will make predictions for these partial decay modes, which are not yet measured.

II. FORMALISM

We follow the steps of Ref. [23] for the weak decay of B mesons leading to the hadronic resonances in the

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final state, generalized to the weak decay of Λ_b baryons into baryonic resonances in Ref. [24]. In Ref. [24] the $\Lambda_b \rightarrow J/\psi K^- p$ and $\Lambda_b \rightarrow J/\psi \pi \Sigma$ reaction in the region of $\Lambda(1405)$ for $K^- p$ ($\pi \Sigma$) were studied, and predictions were made for the $K^- p$ invariant mass distributions, which were confirmed by experiment later in the LHCb work disclosing pentaquark states [25]. The analysis of Ref. [24] also predicted that the $K^- p$ and $\pi \Sigma$ would be produced with isospin $I = 0$, which was also confirmed in Ref. [25] since their partial wave analysis only gave J/ψ and Λ^* states. Work along the same lines as Ref. [24] was done in Ref. [26] in the decay of Λ_c leading to $\Lambda(1405)$ and $\Lambda(1670)$, and in Ref. [27] in the decay of $\Lambda_b \rightarrow J/\psi K \Xi$. The idea of Ref. [24] to the present case proceeds as depicted in Fig. 1.

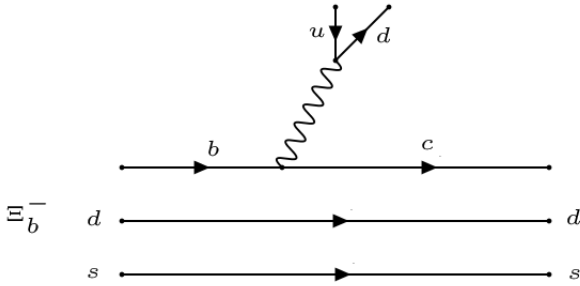


FIG. 1: Diagrammatic representation of the weak decay $\Xi_b^- \rightarrow \pi^- \Xi_c^*$.

The first point to take into account is that in the Ξ_b^- the ds pair has spin $S = 0$. Symmetry of the wave function requires the flavour combination $ds - sd$, and color provides the antisymmetry. The next step is the hadronization of the final $c ds$ state into meson-baryon pairs.

We must consider some basic facts:

1. The ds quarks are spectators in the process. They have $S = 0$ and come in the combination $\frac{1}{\sqrt{2}}(ds - sd)$.
2. The final Ξ_c^* resonances have negative parity, and they will be generated from the meson-baryon interaction in S -wave. Since the pair ds has positive

parity, it is the c quark that must carry the negative parity and hence it will be produced in p -wave ($L = 1$) in the weak interaction of Fig. 1.

3. The c quark will be incorporated into a final $D(D^*)$ meson and thus will go back to its ground state. Hence, the hadronization, introducing $(\bar{u}u + \bar{d}d + \bar{s}s)$ with the quantum numbers of the vacuum, must involve the c quark.

With these constraints the hadronization proceeds as shown in Fig. 2.

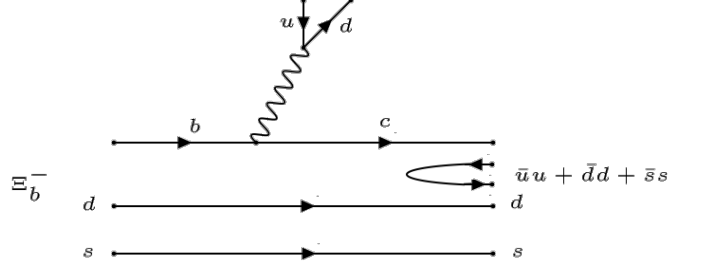


FIG. 2: Hadronization after the weak process in Fig. 1 to produce a meson-baryon pair in the final state.

Technically the hadronization is implemented as follows: The Ξ_b^- state has a flavour function

$$|\Xi_b^-\rangle \equiv \frac{1}{\sqrt{2}} |b(ds - sd)\rangle, \quad (1)$$

and after the weak decay the b quark is substituted by a c quark and we have a state

$$|H\rangle = \frac{1}{\sqrt{2}} |c(ds - sd)\rangle. \quad (2)$$

With the hadronization we have now

$$|H'\rangle = \frac{1}{\sqrt{2}} \sum_{i=1}^3 |P_{4i} q_i (ds - sd)\rangle, \quad (3)$$

where P_{ij} are the $q\bar{q}$ matrix elements.

Next we write the $q\bar{q}$ matrix in terms of the physical mesons, $P \rightarrow \phi$, with ϕ given by

$$\phi \equiv \begin{pmatrix} \frac{1}{\sqrt{2}}\pi^0 + \frac{1}{\sqrt{3}}\eta + \frac{1}{\sqrt{6}}\eta' & \pi^+ & K^+ & \bar{D}^0 \\ \pi^- & -\frac{1}{\sqrt{2}}\pi^0 + \frac{1}{\sqrt{3}}\eta + \frac{1}{\sqrt{6}}\eta' & K^0 & D^- \\ K^- & \bar{K}^0 & -\frac{1}{\sqrt{3}}\eta + \sqrt{\frac{2}{3}}\eta' & D_s^- \\ D^0 & D^+ & D_s^+ & \eta_c \end{pmatrix}. \quad (4)$$

Then we can write

$$|H'\rangle = \frac{1}{\sqrt{2}} [|D^0 u(ds - sd)\rangle + |D^+ d(ds - sd)\rangle + |D_s^+ s(ds - sd)\rangle]. \quad (5)$$

The last state in Eq. (5) contains two extra s quarks and corresponds to a more massive component that we omit in our study.

Next we see that we have a mixed antisymmetric component for the baryonic states of three quarks. If we evaluate the overlap with the mixed antisymmetric representations of the Σ^- , Σ^0 , Λ^0 states [28], we find

$$|H'\rangle = \frac{1}{\sqrt{2}} |D^0 \Sigma^0\rangle + |D^+ \Sigma^-\rangle - \frac{1}{\sqrt{6}} |D^0 \Lambda\rangle. \quad (6)$$

Yet, we have to be careful here with the phase conventions. By looking at the phase convention of Ref. [28] and the one inherent in the baryon octet matrix,

$$B = \begin{pmatrix} \frac{1}{\sqrt{2}} \Sigma^0 + \frac{1}{\sqrt{6}} \Lambda & \Sigma^+ & p \\ \Sigma^- & -\frac{1}{\sqrt{2}} \Sigma^0 + \frac{1}{\sqrt{6}} \Lambda & n \\ \Xi^- & \Xi^0 & -\frac{2}{\sqrt{6}} \Lambda \end{pmatrix}, \quad (7)$$

which is used in the chiral Lagrangians, one can see that one must change the phases of Σ^+ , Λ , Ξ^0 from [28] to agree with the chiral Lagrangians¹.

With this clarification about the phases, the state that we obtain consistent with the chiral convention is:

$$|H'\rangle = \frac{1}{\sqrt{2}} |D^0 \Sigma^0\rangle + |D^+ \Sigma^-\rangle + \frac{1}{\sqrt{6}} |D^0 \Lambda\rangle. \quad (8)$$

We also mention the phase convention for mesons in terms of isospin states, where $|\pi^+\rangle = -|1, 1\rangle$, $|K^-\rangle = -|\frac{1}{2}, -\frac{1}{2}\rangle$, $|D^0\rangle = -|\frac{1}{2}, -\frac{1}{2}\rangle$, and for baryons $\Sigma^+ = -|1, 1\rangle$, $\Xi^- = -|\frac{1}{2}, -\frac{1}{2}\rangle$.

In terms of isospin, $|H'\rangle$ can be written as

$$|H'\rangle = -\sqrt{\frac{3}{2}} \left| \Sigma D(J = \frac{1}{2}) \right\rangle + \frac{1}{\sqrt{6}} \left| \Lambda D(J = \frac{1}{2}) \right\rangle. \quad (9)$$

For D^* production the flavour counting is the same and we would have the same combination substituting D by D^* .

¹ One way to see this is to take the singlet baryon state of Ref. [28] with a minus sign, introduce the hadronization with $\bar{u}u + \bar{d}d + \bar{s}s$ as we have done before and see the meson-baryon content. Then we compare this result with $\text{Tr}(B \cdot \phi)$ obtained with the octet of mesons Eq. (4) for ϕ (taking only the 3×3 part of the matrix), and Eq. (7) for B . The matrix ϕ contains also a singlet of mesons, the octet matrix is the same putting in the diagonal $(\frac{\pi^0}{\sqrt{2}} + \frac{\eta_8}{\sqrt{6}}, -\frac{\pi^0}{\sqrt{2}} + \frac{\eta_8}{\sqrt{6}}, -\frac{2\eta_8}{\sqrt{6}})$.

A. The weak vertex

One must evaluate the weak transition matrix elements. For this we follow the approach in Ref. [21]. The vertex $W^- \rightarrow \pi^-$ is of the type [29, 30]

$$\mathcal{L}_{W\pi} \sim W^\mu \partial_\mu \phi, \quad (10)$$

while the bcW vertex is of the type

$$\mathcal{L}_{qWq} \propto \bar{q}_{\text{fin}} W_\mu \gamma^\mu (1 - \gamma_5) q_{\text{in}}. \quad (11)$$

Since we are dealing with heavy quarks, as in [21] we keep the dominant terms in a nonrelativistic expansion, the γ^0 , $\gamma^i \gamma^5$ ($i = 1, 2, 3$) and, hence, combining the two former Lagrangians we obtain a structure for the weak transition at the quark level

$$V_P \sim q^0 + \vec{\sigma} \cdot \vec{q}, \quad (12)$$

with q^μ the four-momentum of the pion.

In Ref. [21] the operator in Eq. (12), which acts at the quark level between the b and c quarks, was converted into an operator acting over the Λ_c^* and Λ_b at the macroscopical level with the result

$$V_P \sim \left\{ \left(i \frac{q^0}{q} \vec{\sigma} \cdot \vec{q} + iq \right) \delta_{J, \frac{1}{2}} - i \frac{q^0}{q} \sqrt{3} \vec{S}^+ \cdot \vec{q} \delta_{J, \frac{3}{2}} \right\} \text{ME}(q), \quad (13)$$

where \vec{S}^+ is the spin transition operator from spin $\frac{1}{2}$ to spin $\frac{3}{2}$ normalized such that

$$\langle M' | S_\mu^+ | M \rangle = \mathcal{C}(\frac{1}{2}, 1, \frac{3}{2}; M, \mu, M'), \quad (14)$$

with μ in the spherical basis and $\mathcal{C}(\frac{1}{2}, 1, \frac{3}{2}; M, \mu, M')$ the Clebsch-Gordan coefficients. $\text{ME}(q)$ is the quark matrix element for the radial wave functions

$$\text{ME}(q) = \int dr r^2 j_1(qr) \phi_{\text{in}}(r) \phi_{\text{fin}}^*(r). \quad (15)$$

Here we do the same and the macroscopic states are Ξ_c^* and Ξ_b respectively.

Since we require ratios of production rates, the matrix element $\text{ME}(q)$ cancels in the ratio and what matters to differentiate the cases with spin $\frac{1}{2}$ and $\frac{3}{2}$ is the operator in Eq. (13). One should note that the presence of the factor $j_1(qr)$ in Eq. (15) is due to the fact that the c quark is created with $L = 1$ as we discussed previously.

B. The spin structure in the hadronization

The next issue is to see how the hadronization affects the cases of DB or D^*B (with $B = \Sigma, \Lambda$) production in

spin $J = \frac{1}{2}$ or $\frac{3}{2}$. For this we follow again the approach of Ref. [21]. The calculation proceeds as follows:

1. The $\bar{q}q$ pair is created with $J = 0^+$. Since the \bar{q} has negative intrinsic parity we need $L = 1$ in the quarks to restore the positive parity and this forces the $\bar{q}q$ pair to come with spin $S = 1$ to give $J = 0$. This is the essence of the 3P_0 model [28, 31].
2. Since what we want is to elaborate the spin dependence of the matrix elements, we assume a zero range interaction, as is also done in similar problems like the study of pairing in nuclei [32, 33].
3. Since the d , s quarks are spectators and carry $J = 0$, the total angular momentum of the Ξ_c^* is the

same as the angular momentum of the c quark after the weak production.

4. The angular momentum of the c quark and the $\bar{q}q$ pair are recombined to give $L' = 0$, since all quarks are in their ground state in the $D\Sigma$, $D^*\Sigma$, $D\Lambda$, and $D^*\Lambda$ final states. The total angular momentum of the c quark and that of the \bar{q} of the $\bar{q}q$ pair are recombined to give $j = 0, 1$, for the D or D^* production. The total angular momentum of the q from the $\bar{q}q$ pair determines the spin of the baryon Ξ_c^* since the ds quarks carry spin zero. The Clebsch-Gordan coefficient appearing in the different combinations are recombined to give a Racah coefficient [34] and the final result is

$$|J, M; c\rangle |0, 0; \bar{q}q\rangle |0, 0; cd\rangle = \sum_j \mathcal{C}(j, J) |J, M; \text{meson-baryon}\rangle, \quad (16)$$

where the coefficients $\mathcal{C}(j, J)$ are given in Table I.

What we have done so far is to obtain the angular structure of the mechanism for $DB(D^*B)$ production, but we finally want to have the production of the resonances $\Xi_c^0(2790)$ and $\Xi_c^0(2815)$. The way to produce these dynamically generated resonances is depicted in Fig. 3. It involves the amplitudes for $\Xi_b \rightarrow \pi^- D(D^*)B$ production studied before, together with the $D(D^*)B$ loop functions and the couplings of the Ξ_c^* resonance to these

meson-baryon components.

The width for $\Xi_b \rightarrow \pi^- \Xi_c^*$ is given by

$$\Gamma_{\Xi_b \rightarrow \pi^- \Xi_c^*} = \frac{1}{2\pi} \frac{M_{\Xi_c^*}}{M_{\Xi_b}} q \sum \sum |t|^2, \quad (17)$$

with q the momentum of the pion in this decay.

By combining Eqs. (9), (13), (16), we obtain

$$J = \frac{1}{2} : \sum \sum |t|^2 = C^2 (q^2 + \omega_\pi^2) \left| \frac{1}{2} \left(-\sqrt{\frac{3}{2}} \right) G_{\Sigma D} g_{R, \Sigma D} + \frac{1}{2} \frac{1}{\sqrt{6}} G_{\Lambda D} g_{R, \Lambda D} \right. \\ \left. + \frac{1}{2\sqrt{3}} \left(-\sqrt{\frac{3}{2}} \right) G_{\Sigma D^*} g_{R, \Sigma D^*} + \frac{1}{2\sqrt{3}} \frac{1}{\sqrt{6}} G_{\Lambda D^*} g_{R, \Lambda D^*} \right|^2, \quad (18)$$

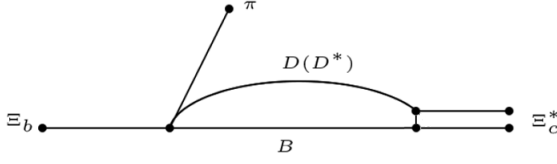
and

$$J = \frac{3}{2} : \sum \sum |t|^2 = C^2 2\omega_\pi^2 \left| \frac{1}{\sqrt{3}} \left(-\sqrt{\frac{3}{2}} \right) G_{\Sigma D^*} g_{R, \Sigma D^*} + \frac{1}{\sqrt{3}} \frac{1}{\sqrt{6}} G_{\Lambda D^*} g_{R, \Lambda D^*} \right|^2, \quad (19)$$

where ω_π is the pion energy $\sqrt{m_\pi^2 + q^2}$, and G_{BD} , G_{BD^*} are the loop functions for the propagator of $BD(BD^*)$ in the resonance formation mechanism of Fig. 3, and $g_{R, BD(BD^*)}$ the coupling of the resonance Ξ_c^* to any of the states $BD(BD^*)$. C in Eqs. (18) (19) is a factor that

contains the matrix element $\text{ME}(q)$ and constants of the weak interaction. Since the mass of the two Ξ_c^* that we investigate are not very different, then we assume C to be a constant that cancels in the ratio of the rates for the

$\mathcal{C}(j, J)$	$J = \frac{1}{2}$	$J = \frac{3}{2}$
(pseudoscalar) $j = 0$	$\frac{1}{4\pi} \frac{1}{2}$	0
(vector) $j = 1$	$\frac{1}{4\pi} \frac{1}{2\sqrt{3}}$	$-\frac{1}{4\pi} \frac{1}{\sqrt{3}}$

TABLE I: $\mathcal{C}(j, J)$ coefficients in Eq. (16).FIG. 3: Mechanism for the production of the Ξ_c^* resonances by production of $D(D^*) \Sigma(\Lambda)$ and coupling of the meson-baryon components to Ξ_c^* .

production of the two resonances. In this case we find

$$R_1 \equiv \frac{\Gamma_{\Xi_b \rightarrow \pi^- \Xi_c(1)}}{\Gamma_{\Xi_b \rightarrow \pi^- \Xi_c(2)}} = \frac{M_{\Xi_c(1)} p_\pi(1) \overline{\sum} \sum |t|^2(1)}{M_{\Xi_c(2)} p_\pi(2) \overline{\sum} \sum |t|^2(2)}, \quad (20)$$

where 1,2 refer to the $\Xi_c(2790)$ and $\Xi_c(2815)$ respectively.

The case of D_s^- production instead of π^- is identical. Instead of $\bar{u}d$ coupling to W we now have $\bar{c}s$, which is equally Cabbibo favoured and is proportional to $\cos \theta_C$ in both cases. The only difference in this case is that the momentum of the D_s^- is smaller than in the case of pion in production. The momenta of D_s^- in the cases $\Xi_c(2790)$ and $\Xi_c(2815)$ are very similar and, hence, by analogy to Eq. (20) we can write

$$R_2 \equiv \frac{\Gamma_{\Xi_b \rightarrow D_s^- \Xi_c(1)}}{\Gamma_{\Xi_b \rightarrow D_s^- \Xi_c(2)}} = \frac{M_{\Xi_c(1)} p_{D_s^-}(1) \overline{\sum} \sum |t|^2(1)}{M_{\Xi_c(2)} p_{D_s^-}(2) \overline{\sum} \sum |t|^2(2)}, \quad (21)$$

with $p_{D_s^-}(1,2)$ evaluated for the $\Xi_c(2790)$ and $\Xi_c(2815)$ respectively, and $\overline{\sum} \sum |t|^2(1,2)$ have to be reevaluated with the new momentum.

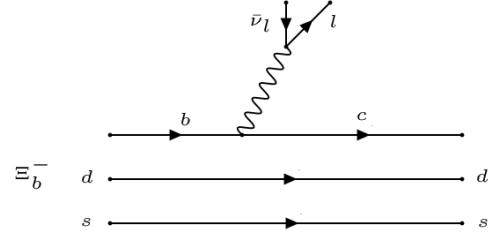
If we assume that $\text{ME}(q)$ is not very different in the case of π^- or D_s^- production we can also write

$$R_3 \equiv \frac{\Gamma_{\Xi_b \rightarrow D_s^- \Xi_c(1)}}{\Gamma_{\Xi_b \rightarrow \pi^- \Xi_c(1)}} = \frac{p_{D_s^-}(1) \overline{\sum} \sum |t|^2(1, D_s^-)}{p_{\pi^-}(1) \overline{\sum} \sum |t|^2(1, \pi^-)}. \quad (22)$$

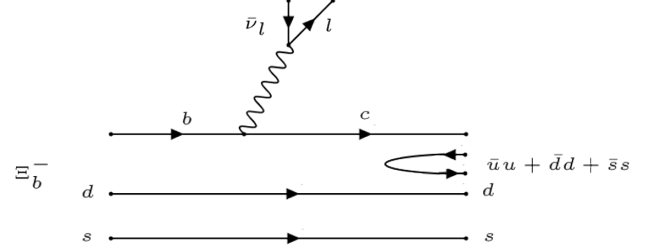
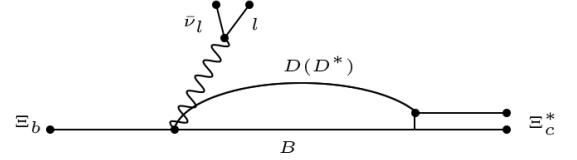
We expect this equation to hold only at the qualitative level since $\text{ME}(q)$ is not necessarily the same for these two different values of q .

III. SEMILEPTONIC DECAY

The semileptonic process, $\Xi_b \rightarrow \bar{\nu}_l l \Xi_c^0(2790)$, $\Xi_b \rightarrow \bar{\nu}_l l \Xi_c^0(2815)$ proceeds in a similar way but instead of a π^-



(a) First step at quark level.

(b) Hadronization to produce $D(D^*) B$.(c) Propagation of $D(D^*) B$ and coupling to the Ξ_c^* .FIG. 4: Different steps of Ξ_c^* production in the $\Xi_b \rightarrow \bar{\nu}_l l \Xi_c^*$ process.

we have $\bar{\nu}_l l$ production. The semileptonic decays of BD hadrons along the lines described here have been studied in Refs. [35, 36]. The weak decay of $\Lambda_c \rightarrow \bar{\nu}_l l \Lambda(1405)$ is addressed in Ref. [37] and the $\Lambda_b \rightarrow \bar{\nu}_l l \Lambda_c(2595)$ and $\Lambda_b \rightarrow \bar{\nu}_l l \Lambda_c(2625)$ in Ref. [22]. The first step from the $\Xi_b \rightarrow \bar{\nu}_l l \Xi_c^*$ is shown in Fig. 4a

The only difference with the nonleptonic decay studied in the former sections is the coupling of W to $\bar{\nu}_l l$. Following Ref. [35] we have, for the combined $W \bar{\nu}_l l$ and Wcb vertices,

$$t = -iG_F \frac{V_{bc}}{\sqrt{2}} L^\alpha Q_\alpha, \quad (23)$$

with G_F the Fermi coupling constant, V_{bc} the Cabbibo-Kobayashi-Maskawa matrix element for the $b \rightarrow c$ tran-

sition, and L^α , Q_α the leptonic and quark currents:

$$L^\alpha = \bar{u}_l \gamma^\alpha (1 - \gamma_5) u_{\nu_l}, \quad (24a)$$

$$Q_\alpha = \bar{u}_c \gamma_\alpha (1 - \gamma_5) u_b. \quad (24b)$$

Once again we retain γ^0 and $\gamma^i \gamma_5$ from the quark matrix elements, which are the leading terms in a nonrelativistic reduction. Actually the $\bar{\nu}_l l$ pair comes out with a large momentum [22] and the momenta of the baryons are small.

The first step in Fig. 4a produces a different structure from Eq. (12) in the nonleptonic case, and one finds [22]

$$\sum_{\text{lepton pol.}} L^\alpha L^\dagger{}^\beta Q_\alpha Q_\beta^\dagger = \frac{8}{m_\nu m_l} p_\nu p_l. \quad (25)$$

The rest of the work needed is identical to the one in the nonleptonic case of the former sections. One can also do an angle integration analytically in the evaluation of Γ and one finally obtains

$$\frac{d\Gamma}{dM_{\text{inv}}(\bar{\nu}_l l)} = \frac{M_{\Xi_c^*}}{M_{\Xi_b}} 2m_\nu 2m_l \frac{1}{(2\pi)^3} p_{\Xi_c^*} \tilde{p}_l \overline{\sum \sum} |t|^2, \quad (26)$$

where $p_{\Xi_c^*}$ is the Ξ_c^* momentum in the Ξ_b rest frame and \tilde{p}_l the lepton momentum in the $\nu_l l$ rest frame, and $\overline{\sum \sum} |t|^2$ is given by [22]

$$\overline{\sum \sum} |t|^2 = C'^2 \frac{8}{m_\nu m_l} \frac{1}{M_{\Xi_b}^2} \left(\frac{M_{\text{inv}}}{2} \right)^2 \left[\tilde{E}_{\Xi_b}^2 - \frac{1}{3} \tilde{p}_{\Xi_b}^2 \right] A_J V_{\text{had}}(J), \quad (27)$$

with

$$J = \frac{1}{2} : A_J V_{\text{had}}(J) = \left| \frac{1}{2} \left(-\sqrt{\frac{3}{2}} \right) G_{\Sigma D} g_{R, \Sigma D} + \frac{1}{2} \frac{1}{\sqrt{6}} G_{\Lambda D} g_{R, \Lambda D} + \right. \\ \left. + \frac{1}{2\sqrt{3}} \left(-\sqrt{\frac{3}{2}} \right) G_{\Sigma D^*} g_{R, \Sigma D^*} + \frac{1}{2\sqrt{3}} \frac{1}{\sqrt{6}} G_{\Lambda D^*} g_{R, \Lambda D^*} \right|^2, \quad (28)$$

and,

$$J = \frac{3}{2} : A_J V_{\text{had}}(J) = 2 \left| \frac{1}{\sqrt{3}} \left(-\sqrt{\frac{3}{2}} \right) G_{\Sigma D^*} g_{R, \Sigma D^*} + \frac{1}{\sqrt{3}} \frac{1}{\sqrt{6}} G_{\Lambda D^*} g_{R, \Lambda D^*} \right|^2, \quad (29)$$

where G_{BD} , G_{BD^*} and $g_{R, BD}$, g_{R, BD^*} are the same as in the nonleptonic decay and C' is again a factor that contains the matrix element $\text{ME}(q)$ evaluated at the proper value of q . A novelty here is that q is not constant when one integrates $\frac{d\Gamma}{dM_{\text{inv}}}$ over M_{inv} . However, the fact that M_{inv} peaks around the maximum allowed in the Dalitz plot [22], as we show in Fig. 5 for the present case, allows us to consider C' constant over the whole range of M_{inv} .

The magnitudes \tilde{E}_{Ξ_b} and \tilde{p}_{Ξ_b} in Eq. (26) are the energies of Ξ_b and its momentum in the rest frame of the $\nu_l l$ pair which are given by [35]

$$\tilde{E}_{\Xi_b} = \frac{M_{\Xi_b}^2 + M_{\text{inv}}^2 - M_{\Xi_c^*}^2}{2M_{\text{inv}}}, \quad (30a)$$

$$\tilde{p}_{\Xi_b} = \frac{\lambda^{\frac{1}{2}}(M_{\Xi_b}^2, M_{\text{inv}}^2, M_{\Xi_c^*}^2)}{2M_{\text{inv}}}, \quad (30b)$$

with $\lambda(x, y, z)$ the ordinary Källén function.

An approximate value for the ratio of the semileptonic production for the two resonances is given by

$$R = \frac{\Gamma_{\Xi_b \rightarrow \bar{\nu}_l l} \Xi_c(2790)}{\Gamma_{\Xi_b \rightarrow \bar{\nu}_l l} \Xi_c(2815)} = \frac{A_{\frac{1}{2}} V_{\text{had}}\left(\frac{1}{2}\right)}{A_{\frac{3}{2}} V_{\text{had}}\left(\frac{3}{2}\right)}. \quad (31)$$

IV. RESULTS

We use the values of $g_{R, \Sigma D}$, $g_{R, \Sigma D^*}$, $g_{R, \Lambda D}$, $g_{R, \Lambda D^*}$ and of the $G_{\Sigma D}$, $G_{\Sigma D^*}$, $G_{\Lambda D}$, $G_{\Lambda D^*}$ from Ref. [18] which we have redone in order to evaluate the complex couplings and the G functions since only the modulus of $g_{R, i}$ were given there and the values of G_i were not tabulated. We give all this information in Tables II and III. The G functions are taken from dimensional regularization subtracting the value of G at $s = \alpha (M_{\text{th}}^2 + m_{\text{th}}^2)$

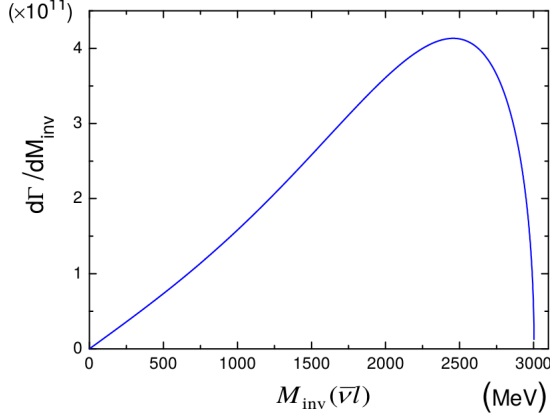


FIG. 5: The invariant mass distribution for $\bar{\nu}_l l$ in the $\Xi_b \rightarrow \bar{\nu}_l l \Xi_c(2790)$. The one for the $\Xi_b \rightarrow \bar{\nu}_l l \Xi_c(2815)$ decay is very similar.

	ΣD	ΣD^*	ΛD	ΛD^*
g	$1.209+i0.107$	$-0.801-i0.286$	$1.438-i0.864$	$-0.602+i0.584$
Gg	$-6.734-i0.270$	$3.483+i1.073$	$-8.570+i5.777$	$2.609-i2.777$

TABLE II: Values of g and Gg for the different channels for the resonance $\Xi_c^0(2790)(\frac{1}{2}^-)$.

with $\alpha = 0.9698$ and $M_{\text{th}} + m_{\text{th}}$ the mass of the lightest hadronic channel of all the coupled channels for a given quantum number [38].

Using the values in Tables II and III and Eq. (20) we obtain

$$R_1 = \frac{\Gamma_{\Xi_b \rightarrow \pi^- \Xi_c^0(2790)}}{\Gamma_{\Xi_b \rightarrow \pi^- \Xi_c^0(2815)}} = 0.404, \quad (32)$$

and for Eq. (21)

$$R_2 = \frac{\Gamma_{\Xi_b \rightarrow D_s^- \Xi_c^0(2790)}}{\Gamma_{\Xi_b \rightarrow D_s^- \Xi_c^0(2815)}} = 0.287. \quad (33)$$

Similarly we can obtain for Eq. (22)

$$R_3 = \frac{\Gamma_{\Xi_b \rightarrow D_s^- \Xi_c^0(2790)}}{\Gamma_{\Xi_b \rightarrow \pi^- \Xi_c^0(2790)}} = 0.686. \quad (34)$$

In order to see how sensitive these rates are to the values of the D^*B couplings we reevaluate them by first setting them to zero then changing their sign. The results we obtain are shown in Table IV.

	ΛD^*	ΣD^*
g	$2.344-i0.605$	$0.783+i0.483$
Gg	$-12.286+i4.237$	$-4.108-i2.123$

TABLE III: Values of g and Gg for the different channels for the resonance $\Xi_c^0(2815)(\frac{3}{2}^-)$.

	R_1	R_2	R_3
$g_{R,\Sigma D^*} = 0$	0.882	0.626	0.686
$g_{R,\Lambda D^*} = 0$	0.21	0.149	0.686
$g_{R,\Sigma D^*} \rightarrow -g_{R,\Sigma D^*}$	0.512	0.363	0.686
$g_{R,\Lambda D^*} \rightarrow -g_{R,\Lambda D^*}$	0.07	0.05	0.686

TABLE IV: Values of R_1 , R_2 , R_3 by both changing the sign of the g_{R,BD^*} couplings and setting them to zero.

	R
$g_{R,\Sigma D^*} = 0$	0.452
$g_{R,\Lambda D^*} = 0$	0.108
$g_{R,\Sigma D^*} \rightarrow -g_{R,\Sigma D^*}$	0.263
$g_{R,\Lambda D^*} \rightarrow -g_{R,\Lambda D^*}$	0.036

TABLE V: Values of R of the semileptonic decay, obtained by both changing the sign of the g_{R,BD^*} couplings and setting them to zero.

As we can see, the results shown above tell us the relevance of the D^*B components in the build up of these resonances.

As for the sector of the semileptonic decay rates corresponding to Eq. (31) we find that

$$R = \frac{\Gamma_{\Xi_b \rightarrow \bar{\nu}_l l \Xi_c(2790)}}{\Gamma_{\Xi_b \rightarrow \bar{\nu}_l l \Xi_c(2815)}} = 0.201, \quad (35)$$

and if we integrate Eq. (26) we find

$$R = 0.207. \quad (36)$$

As we can see, the numbers are essentially the same. Once again, if the couplings to D^*B states are changes we obtain different results, shown in Table V.

V. CONCLUSION

We have studied the nonleptonic $\Xi_b^- \rightarrow M + \Xi_c^*$, with $M = \pi^-$, D_s^- and $\Xi_c^* = \Xi_c^0(2790)(\frac{1}{2}^-)$, $\Xi_c^0(2815)(\frac{3}{2}^-)$. We have assumed that the Ξ_c^* resonances are dynamically generated from the MB and VB interaction, as done in Ref. [18]. We saw that the present decays only involved the $D\Lambda$, $D\Sigma$, $D^*\Lambda$, $D^*\Sigma$ channels and we took the needed couplings from that work. Given the fact that the momentum of the meson M is very similar for the case of the production of the two resonances (since their masses are very close) we could eliminate in the ratio of widths the matrix element at the quark level involving the wave functions of the b and c quarks. Then, only factors related to the spin structure of the channels and the couplings of the hadronic model for the resonances were relevant, which tells us that the measurement of these partial decay widths are relevant to learn details on the nature of the Ξ_c^* resonances. With more uncertainty we

were able to also predict the ratio of $\Xi_b^- \rightarrow \pi^- \Xi_c^*$ and $\Xi_b^- \rightarrow D_s^- \Xi_c^*$ for the same resonance.

We also evaluated in the semileptonic rates. In this case we can only evaluate one ratio, the semileptonic decay $\Xi_b^- \rightarrow \bar{\nu}_l l \Xi_c^*$ for the $\Xi_c^0(2790)$ and $\Xi_c^0(2815)$ resonances. Once again, the predictions will be valuable when these partial decay width can be measured. We should stress that both the nonleptonic and semileptonic decay widths are measured for the case of $\Lambda_b \rightarrow \pi^- \Lambda_c(2595)$, $\Lambda_b \rightarrow \pi^- \Lambda_c(2625)$ and $\Lambda_b \rightarrow \bar{\nu}_l l \Lambda_c(2595)$ and $\Lambda_b \rightarrow \bar{\nu}_l l \Lambda_c(2625)$ and the method used here gave results in agreement with experiment [21, 22], so we are confident that the predictions done here are fair. In any case the experimental result could test the accuracy of the model of Ref. [18], which is one of the possible ways to address the molecular states, with a particular dynamics based on a particular symmetry.

We also checked that the results were sensitive to the couplings of the D^*B components and confirmation of this feature by experiment could give a boost to the rel-

evance of the mixing of pseudoscalar-baryon and vector-baryon components in the building up of the molecular baryonic states, a subject which is catching up in the hadronic community [15–18, 39–41].

Acknowledgments

R. P. Pavao wishes to thank the Generalitat Valenciana in the program Santiago Grisolia. This work is partly supported by the National Natural Science Foundation of China under Grants No. 11565007, 11647309 and No. 11547307. This work is also partly supported by the Spanish Ministerio de Economía y Competitividad and European FEDER funds under the contract number FIS2011-28853-C02-01, FIS2011-28853-C02-02, FIS2014-57026-REDT, FIS2014-51948-C2-1-P, and FIS2014-51948-C2-2-P, and the Generalitat Valenciana in the program Prometeo II-2014/068.

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